

FINITE BRST-ANTIBRST TRANSFORMATIONS FOR THE THEORIES WITH GAUGE GROUP.

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Following our recent study [P.Yu. Moshin, A.A. Reshetnyak, Nucl. Phys. B 888 (2014) 92], we discuss the notion of finite BRST-antiBRST transformations, with a doublet λ_a , $a = 1, 2$, of anticommuting (both global and field-dependent) Grassmann parameters. It turns out that the global finite BRST-antiBRST transformations form a 2-parametric Abelian supergroup. We find an explicit Jacobian corresponding to this change of variables in the theories with a gauge group. Special field-dependent BRST-antiBRST transformations for the Yang–Mills path integral with s_a -potential (functionally-dependent) parameters $\lambda_a = s_a \Lambda$ generated by a finite even-valued functional Λ and the anticommuting generators s_a of BRST-antiBRST transformations, amount to a precise change of the gauge-fixing functional. This proves the independence of the vacuum functional under such BRST-antiBRST transformations and leads to the presence of modified Ward identities. The form of transformation parameters that induces a change of the gauge in the path integral is found and is exactly evaluated for connecting two arbitrary R_ξ -like gauges. The finite field-dependent BRST-antiBRST transformations are used to generalize the Gribov horizon functional h_0 , in the Landau gauge of the BRST-antiBRST setting in the Gribov–Zwanziger model, and to find h_ξ corresponding to general R_ξ -like gauges in the form compatible with a gauge-independent S -matrix.

Keywords: *gauge theories, BRST-antiBRST Lagrangian quantization, Yang–Mills theory, Gribov–Zwanziger model, field-dependent BRST-antiBRST transformations*

1 Motivations

The special supersymmetries known as BRST symmetry [1, 2] and BRST-antiBRST symmetry [3, 4, 5] provide a basis of the modern quantization methods for gauge theories [6, 7]. These symmetries feature the presence of a Grassmann-odd parameter μ and two Grassmann-odd parameters $(\mu, \bar{\mu})$, respectively. The latter parameters in the $\text{Sp}(2)$ -covariant schemes of generalized Hamiltonian [8] and Lagrangian [9, 10] quantization (see [11] as well) $(\mu, \bar{\mu}) \equiv (\mu_1, \mu_2) = \mu_a$ form an $\text{Sp}(2)$ -doublet. These infinitesimal odd-valued parameters may be regarded both as constants and field-dependent functionals, used, respectively, to derive the Ward identities and to establish the gauge-independence of the partition function in the path integral approach.

In [12], the BRST transformations with a finite field-dependent parameter (FFDBRST) in the gauge theories with a gauge group, i.e., Yang–Mills (YM) theories, whose quantum action is constructed by the Faddeev–Popov (FP) rules [13], were first introduced by means of a functional equation for the parameter in question and used to provide the path integral with such a change of variables that would allow one to re-

late the quantum action given in a certain gauge with the one given in a different gauge, however, without solving it in the general setting (see [14] as well). The problem of establishing a relation of the FP action in a certain gauge with the action in a different gauge by using a change of variables induced by a FFDBRST transformation was solved in [15], thereby providing an exact relation between a finite parameter and a finite change of the gauge-fixing condition in terms of the gauge Fermion. In particular, this result implies the conservation of the number of physical degrees of freedom in a given YM theory with respect to FFDBRST transformations, which means the impossibility of relating the Yang–Mills theory to another theory, whose action may contain, in addition to the FP action, certain non-BRST invariant terms (such as the Gribov horizon functional [23] in the Gribov–Zwanziger theory [24], as a consequence of Singer’s result [25]) in the same configuration space.¹

Notice that the solution of a similar problem for arbitrary constrained dynamical systems in the generalized Hamiltonian formalism [16, 17] has been recently proposed in [18], whereas for general gauge theories (possessing reducible gauge symmetries and/or an open gauge algebra) an exact Jacobian generated by FFD-

¹For the study of the Gribov copies in YM theories, using the covariant, Landau, maximal Abelian gauges, see [26, 27, 28, 29, 30].

BRST transformations in the path integral constructed by the BV procedure [19] has been obtained in [20] (see [21] as well), with a solution of the consistency of *soft BRST symmetry breaking* [22].

Recently, we have proposed an extension of BRST-antiBRST transformations to the case of finite (both global and field-dependent) parameters in Yang–Mills [31] and general gauge theories [32, 33, 34], using the Lagrangian and generalized Hamiltonian BRST-antiBRST quantization methods; also see [35]. Here, we review the origin of finite BRST-antiBRST transformations and use their properties to study their influence on the quantum structure of YM theories in the framework of the BRST-antiBRST setting.

We use the conventions introduced in [31]. Unless otherwise specified by an arrow, derivatives with respect to the fields are taken from the right, and those with respect to the corresponding antifields are taken from the left. The raising and lowering of $\text{Sp}(2)$ indices, $s^a = \varepsilon^{ab}s_b$, $s_a = \varepsilon_{ab}s^b$, is carried out by a constant antisymmetric metric tensor ε^{ab} , $\varepsilon^{ac}\varepsilon_{cb} = \delta_b^a$, subject to the normalization $\varepsilon^{12} = 1$.

2 Proposal for Finite Field-Dependent BRST-antiBRST Transformations

The generating functional of Green’s functions for irreducible gauge theories with a closed algebra in the BRST-antiBRST Lagrangian quantization [9, 10] is given by

$$Z_F(J) = \int d\phi \exp \left\{ \frac{i}{\hbar} [S_F(\phi) + J_A \phi^A] \right\} \quad (1)$$

and depends on the sources J_A , with the BRST-antiBRST-invariant quantum action

$$\begin{aligned} S_F(\phi) &= S_0(A) + 1/2 s^a s_a F(A, C) \\ &= S_0(A) + S_{\text{gf}}(A, B) + S_{\text{gh}}(A, C) + S_{\text{add}}(C) \end{aligned} \quad (2)$$

defined in the same total configuration space \mathcal{M} as in the FP method, and parameterized by the respective classical fields, the $\text{Sp}(2)$ -doublet of ghost-antighost fields, and the Nakanishi–Lautrup fields, $\phi^A = (A^i, C^{\alpha a}, B^\alpha)$, with the Grassmann parities $\varepsilon(\phi^A) \equiv \varepsilon_A$, $\varepsilon_A = \varepsilon(A^i, B^\alpha, C^{\alpha a}) \equiv (\varepsilon_i, \varepsilon_\alpha, \varepsilon_\alpha + 1)$, using the condensed notation. The quantities S_0 , F are the respective classical gauge-invariant action and an admissible gauge-fixing Bosonic functional, chosen here in the quadratic approximation for YM theories, with $A^i = A^{\mu n}(x)$ defined in a d -dimensional Minkowski space and taking values in the Lie algebra of the $SU(N)$ gauge group, with $\eta_{\mu\nu} = \text{diag}(-, +, \dots, +)$,

$$S_0 = -1/4 \int d^d x F_{\mu\nu}^n F^{\mu\nu n}, n = 1, \dots, N^2 - 1, \quad (3)$$

in terms of the strength $F^{\mu\nu n} = \partial^{[\mu} A^{\nu]n} + f^{nop} A^{\mu 0} A^{\nu p}$ and

$$F_\xi(A, C) = -\frac{1}{2} \int d^d x \left(A_\mu^m A^{m\mu} - \xi/2 \varepsilon_{ab} C^{ma} C^{mb} \right), \quad (4)$$

which corresponds to the R_ξ -family of gauges with $\chi_\xi(A, B) = \partial_\mu A^{\mu a} + \frac{\xi}{2} B^a = 0$ in the FP rules for YM theories. The remaining terms in (2), i.e., the gauge-fixing term S_{gf} , the ghost term S_{gh} , and the interaction term S_{add} , quartic in C^{ma} , are given by

$$\begin{aligned} S_{\text{gf}} &= \int d^d x \left[(\partial^\mu A_\mu^m) + \xi/2 B^m \right] B^m, \\ S_{\text{gh}} &= \frac{1}{2} \int d^d x (\partial^\mu C^{ma}) D_\mu^{mn} C^{nb} \varepsilon_{ab}, \\ S_{\text{add}} &= -\frac{\xi}{48} \int d^d x f^{mnl} f^{lrs} C^{sa} C^{rc} C^{mb} C^{md} \varepsilon_{ab} \varepsilon_{cd}. \end{aligned} \quad (5)$$

The action (3) is invariant with respect to the infinitesimal gauge transformations $\delta A_\mu^m = D_\mu^{mn} \zeta^n$ with arbitrary functions $\zeta^\alpha \equiv \zeta^n$, $\varepsilon(\alpha) = 0$, defined in $R^{1, d-1}$, whereas the infinitesimal BRST-antiBRST transformations $\delta \phi^A = s^a \phi^A \mu_a$ for YM theories are given in terms of anticommuting generators $s^a : s^a s^a b + s^b s^a = 0$,

$$\begin{aligned} s^a A_\mu^m &= D_\mu^{mn} C^{ma}, \\ s^a \delta B^m &= 1/2 f^{nml} (B^l C^{na} + (1/6) f^{lrs} C^{sb} C^{ra} C^{mc} \varepsilon_{cb}), \\ s^b C^{ma} &= (\varepsilon^{ab} B^m - (1/2) f^{mnl} C^{la} C^{nb}), \end{aligned} \quad (6)$$

leaving the action S_F and the integrand \mathcal{I}_ϕ^F in $Z_F(0) \equiv \int \mathcal{I}_\phi^F$ invariant only in the 1-st order in powers of μ_a .

To restore the total BRST-antiBRST invariance of S_F and \mathcal{I}_ϕ^F to all orders in μ_a , we have introduced [31] finite transformations of ϕ^A with a doublet λ_a of anticommuting parameters, $\lambda_a \lambda_b + \lambda_b \lambda_a = 0$,

$$\phi^A \rightarrow \phi'^A = \phi^A + \Delta \phi^A = \phi'^A(\phi|\lambda) : \phi'(\phi|0) = \phi, \quad (7)$$

as a solution of the functional equation

$$G(\phi') = G(\phi) \quad \text{if} \quad s^a G(\phi) = 0 \quad (8)$$

for any regular functional $G(\phi)$ invariant under infinitesimal BRST-antiBRST transformations. The general solution of (8) allows one to restore *finite BRST-antiBRST transformations* in a unique way $\phi^A \rightarrow \phi'^A$,

$$\phi'^A = \phi^A \left(1 + \overleftarrow{s}^a \lambda_a + \frac{1}{4} \overleftarrow{s}^2 \lambda^2 \right) \equiv \phi^A \exp(\overleftarrow{s}^a \lambda_a), \quad (9)$$

where the set of elements $\{g(\lambda)\} = \{\exp(\overleftarrow{s}^a \lambda_a)\}$ forms an Abelian two-parametric supergroup with odd-valued generating elements λ_a . The prescription (7), (8) for obtaining finite (group) BRST-antiBRST transformations (9) works perfectly well in the case of changing the configuration space \mathcal{M} to another representation space in which the action of superalgebra of s^a

is realized. Thus, the finite BRST-antiBRST transformations have also been constructed in the generalized Hamiltonian [34] and Lagrangian [32] BRST-antiBRST quantization schemes, which is also supported by the Frobenius theorem. The BRST-antiBRST invariance of \mathcal{I}_ϕ^F implies the validity of the relation

$$\mathcal{I}_{\phi g(\lambda)}^F = \mathcal{I}_\phi^F, \quad (10)$$

with allowance for the fact established in [31] that the global finite transformations, corresponding to $\lambda_a = \text{const}$, respect the integration measure.

3 Jacobian of Finite BRST-antiBRST Transformations

As we have already mentioned, the Jacobian of the change of variables corresponding to the global finite transformations is equal to 1:

$$\mathfrak{S}(\phi) = 0 \implies \text{Sdet} \left(\frac{\delta \phi'}{\delta \phi} \right) = 1 \text{ and } d\phi' = d\phi. \quad (11)$$

We have also examined [31] the finite field-dependent transformations in the particular case of functionally-dependent parameters $\lambda_a = \Lambda \overleftarrow{s}_a$, $s^1 \lambda_1 + s^2 \lambda_2 = -s^2 \Lambda$, with a certain even-valued potential, $\Lambda = \Lambda(\phi)$, inspired by infinitesimal field-dependent BRST-antiBRST transformations with the parameters

$$\mu_a = \frac{i}{2\hbar} \varepsilon_{ab} (\Delta F)_{,A} X^{Ab} = \frac{i}{2\hbar} (s_a \Delta F), \quad (12)$$

which respect the gauge independence of the integrand, and therefore also of the vacuum functional $Z_F(0)$, with accuracy up to the terms linear in ΔF : $\mathcal{I}_{\phi g(\mu(\Delta F))}^F = \mathcal{I}_\phi^{F+\Delta F} + o(\Delta F)$. In the case of finite field-dependent transformations with the group elements $g(\Lambda \overleftarrow{s}_a)$, which now form a non-Abelian 2-parametric supergroup, the superdeterminant of a change of variables takes the form

$$\text{Sdet} \left(\frac{\delta(\phi g(\Lambda \overleftarrow{s}_a))}{\delta \phi} \right) = \left[1 - \frac{1}{2} s^2 \Lambda(\phi) \right]^{-2}, \quad (13)$$

$$d\phi' = d\phi \exp \left\{ \frac{i}{\hbar} \left[i\hbar \ln \left(1 - \frac{1}{2} s^2 \Lambda \right)^2 \right] \right\}. \quad (14)$$

4 Compensation Equation for Yang–Mills Theories in Different Gauges

In view of the invariance of the quantum action $S_F(\phi)$ under (9), the change $\phi^A \rightarrow \phi'^A = \phi^A g(\lambda(\phi))$ induces in (1) the transformation of the integrand \mathcal{I}_ϕ^F

$$\mathcal{I}_{\phi g(\lambda(\phi))}^F = d\phi \exp \left\{ \frac{i}{\hbar} \left[S_F(\phi) + i\hbar \ln \left(1 - \frac{1}{2} s^2 \Lambda \right)^2 \right] \right\}. \quad (15)$$

Due to the explicit form of the initial quantum action $S_F = S_0 - (1/2) F \overleftarrow{s}^2$, the BRST-antiBRST-exact contribution $i\hbar \ln(1 + s^a s_a \Lambda/2)^2$ to the action S_F , resulting from the transformation of the integration measure, can be interpreted as a change of the gauge-fixing functional made in the original integrand \mathcal{I}_ϕ^F ,

$$i\hbar \ln(1 + s^a s_a \Lambda/2)^2 = s^a s_a (\Delta F/2) \quad (16)$$

$$\implies \mathcal{I}_{\phi g(\lambda(\phi))}^F = \mathcal{I}_\phi^{F+\Delta F}, \quad (17)$$

for a certain $\Delta F(\phi|\Lambda)$, whose relation to $\Lambda(\phi)$ is established by (16), referred to in [31] as the *compensation equation* for an unknown parameter $\Lambda(\phi)$, which provides the gauge independence of the vacuum functional, $Z_F(0) = Z_{F+\Delta F}(0)$. An explicit solution to (16), which obeys the solvability condition due to the BRST-antiBRST exactness of its both sides, is given, up to BRST-antiBRST-exact terms, by the relations

$$\begin{aligned} \Lambda(\phi|\Delta F) &= 2\Delta F (s^a s_a \Delta F)^{-1} \left[\exp \left(\frac{1}{4i\hbar} s^b s_b \Delta F \right) - 1 \right] \\ &= \frac{1}{2i\hbar} \Delta F \sum_{n=0}^{\infty} \frac{1}{(n+1)!} \left(\frac{1}{4i\hbar} s^a s_a \Delta F \right)^n. \end{aligned} \quad (18)$$

Conversely, having considered the equation (16) for an unknown ΔF with a given Λ , we obtain

$$\Delta F(\phi) = -2i\hbar \Lambda (s^2 \Lambda)^{-1} \ln(1 - s^2 \Lambda/2)^2. \quad (19)$$

Thus, the field-dependent transformations with the parameters $\lambda_a = s_a \Lambda$ amount to a precise change of the gauge-fixing functional. For instance, in order to relate $Z_{F_\xi}(J)$ to $Z_{F_{\xi+\Delta\xi}}(J)$ in the R_ξ -family of gauges, one has to carry out FDBRST-antiBRST transformations with the parameters

$$\begin{aligned} \lambda_a &= \frac{\Delta\xi}{4i\hbar} \varepsilon_{ab} \int d^d x \left(B^n C^{mb} \sum_{n=0}^{\infty} \frac{1}{(n+1)!} \right) \left[\frac{\Delta\xi}{4i\hbar} \int d^d y \right. \\ &\quad \times \left(B^u B^u - \frac{1}{24} f^{uwt} f^{trs} C^{sc} C^{rp} C^{wd} C^{uq} \varepsilon_{cd} \varepsilon_{pq} \right) \Big]. \end{aligned} \quad (20)$$

5 Gauge Dependence Problem and modified Ward identities

In [33], the property (17) leads to a so-called modified Ward identity, depending on field-dependent parameters, $\lambda_a = \Lambda \overleftarrow{s}_a$, and therefore also on a finite change of the gauge, in view of (18),

$$\begin{aligned} &\left\langle \left\{ 1 + \frac{i}{\hbar} J_A \left[X^{Aa} \lambda_a(\Lambda) - \frac{1}{2} Y^A \lambda^2(\Lambda) \right] - \frac{1}{4} \left(\frac{i}{\hbar} \right)^2 \varepsilon_{ab} \right. \right. \\ &\quad \times J_A X^{Aa} J_B X^{Bb} \lambda^2(\Lambda) \Big\} \left(1 - \frac{1}{2} \Lambda \overleftarrow{s}^2 \right)^{-2} \Big\rangle_{F,J} = 1, \end{aligned} \quad (21)$$

for $(s^a \phi^A, s^2 \phi^A) \equiv (X^{Aa}, -2Y^A)$. The property (17) also implies a relation describing the gauge dependence

of $Z_F(J)$ for a finite change of the gauge $F \rightarrow F + \Delta F$:

$$\begin{aligned} \Delta Z_F(J) = Z_F(J) & \left\langle \frac{i}{\hbar} J_A [X^{Aa} \lambda_a (\phi| - \Delta F) \right. \\ & - \frac{1}{2} Y^A \lambda^2 (\phi| - \Delta F)] - (-1)^{\varepsilon_B} \left(\frac{i}{2\hbar} \right)^2 J_B J_A \\ & \left. \times (X^{Aa} X^{Bb}) \varepsilon_{ab} \lambda^2 (\phi| - \Delta F) \right\rangle_{F,J}, \end{aligned} \quad (22)$$

which means the on-shell ($J = 0$) gauge independence of the conventional S -matrix due to equivalence theorem [36]. In (21), (22), the symbol “ $\langle \mathcal{A} \rangle_{F,J}$ ” for a quantity $\mathcal{A} = \mathcal{A}(\phi)$ denotes the source-dependent average expectation value for a gauge-fixing $F(\phi)$

$$\langle \mathcal{A} \rangle_{F,J} = Z_F^{-1}(J) \int d\phi \mathcal{A}(\phi) \exp \left\{ \frac{i}{\hbar} [S_F + J\phi] \right\}, \quad (23)$$

with $\langle 1 \rangle_{F,J} = 1$. For constant λ_a , the relation (21) implies an $\text{Sp}(2)$ -doublet of the usual Ward identities at the first order in λ_a : $J_A \langle X^{Aa} \rangle_{F,J} = 0$ and a derivative identity at the second order in λ_a :

$$\langle J_A [2Y^A + (i/\hbar) \varepsilon_{ab} X^{Aa} J_B X^{Bb}] \rangle_{F,J} = 0. \quad (24)$$

6 Gribov–Zwanziger Theory in BRST-antiBRST Formulation of Landau and Feynman Gauges

Since the gauge-fixing functional F_0 corresponds to the Landau gauge, we introduce the Gribov horizon functional in the same manner as in [24] for the FP procedure in the Euclidian space:

$$\begin{aligned} h(A) = \int d^d x d^d y f^{mrl} A_\mu^r(x) K^{mn-1}(x; y) \\ f^{nsl} A^{\mu s}(y) + d(N^2 - 1), \end{aligned} \quad (25)$$

with γ^2 , $K^{mn}(x; y)$ being the Gribov mass parameter determined from the gap equation and the FP matrix. A proper BRST-antiBRST-non-invariant action in the F_0 reference frame has the form

$$S_h(\phi) = S_{F_0}(\phi) + \gamma^2 h(\phi). \quad (26)$$

We determine the Gribov–Zwanziger theory in any F_ξ gauges (R_ξ -gauges) in a way compatible with the gauge-independence of the generating functional of Green’s functions in F_0 , where Gribov horizon in the gauge F_ξ should be determined as

$$\begin{aligned} h_\xi = h \left(1 + \frac{1}{2i\hbar} (\overleftarrow{s}^a) (\Delta F_\xi \overleftarrow{s}^a) \sum_{n=0}^{\infty} \frac{1}{(n+1)!} \right. \\ \times \left(-\frac{1}{4i\hbar} \Delta F_\xi \overleftarrow{s}^2 \right)^n - \frac{1}{16\hbar^2} (\overleftarrow{s}^2) (\Delta F_\xi)^2 \\ \left. \times \left[\sum_{n=0}^{\infty} \frac{1}{(n+1)!} \left(-\frac{1}{4i\hbar} \Delta F_\xi \overleftarrow{s}^2 \right)^n \right]^2 \right), \end{aligned} \quad (27)$$

where ΔF_ξ is readily determined with account taken of (20); for details, see [31]. The construction of the Gribov horizon functional h_ξ in the gauge F_ξ , starting from $h_0 = h(A)$ in the gauge F_0 , may be considered as a generalization of the result [37] obtained in the BRST setting of the problem. Considering, instead of h_0 , a functional calculated, e.g., in the Coulomb gauge [38], we may use the above recipe to construct the form of h_ξ in any F_ξ -gauges, including the Landau and Feynman gauges, and compare them with those obtained by the non-perturbative Zwanziger’s recipe [24], in order to verify the consistency of Zwanziger’s recipe with the requirement of gauge independence for the corresponding vacuum functionals.

7 Conclusion

We have proposed the concept of finite BRST-antiBRST and FFDBRST-antiBRST transformations for Yang–Mills theories in the $\text{Sp}(2)$ -covariant Lagrangian quantization. The Jacobian of the change of variables generated by FFDBRST-antiBRST transformations with functionally dependent parameters is exactly calculated. It is established that quantum YM actions in different gauges are related to each other by means of FFDBRST-antiBRST transformations with functionally dependent parameters obtained as solutions of the compensation equation. A new Ward identity and the gauge dependence problem for finite changes of the gauge for the generating functional of Green’s functions are derived and studied. The Gribov–Zwanziger theory in a BRST-antiBRST formulation is suggested and the Gribov horizon functional, h , in the Feynman and arbitrary gauges, starting from h_0 given in Landau gauge, is suggested in a way compatible with the gauge independence of the S -matrix. Concluding, note that Gribov’s problem for YM theories in BRST-antiBRST setting may be elaborated with help of composite fields technique following to the results [39] obtained within BRST setting of the problem.

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